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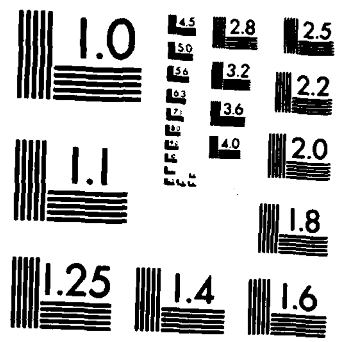

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On the Linear Stability of Two-Dimensional Barium Clouds
I. The Inviscid Case

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Plasma Physics Division

April 30, 1984

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We examine the linear stability of the steepened backside of two-dimensional ionospheric barium clouds. We derive expressions for the growth rate of perturbations on these backsides which are functions of (1) the ratio M of the integrated Pedersen conductivity inside the barium cloud to that of the background ionosphere; (2) the overall shape of the cloud; (3) the plasma density gradient scale length L of the steepened backside; and (4) the perturbation wavenumber k. For sufficiently small kL, the growth time as a function of M traces out the "U-shaped curve" observed in previous studies.

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ON THE LINEAR STABILITY OF TWO-DIMENSIONAL BARIUM CLOUDS
I. THE INVISCID CASE

CONTENTS

ON THE LINEAR STABILITY OF TWO-DIMENSIONAL BARIUM CLOUDS
I. THE INVISCID CASE

I. Introduction

The study of barium cloud phenomenology has been of interest to ionospheric physicists for more than two decades. The release of barium in the ionosphere, followed by its rapid ionization, leads to a complex interaction between the barium plasma and the background ionospheric plasma in the presence of ambient electric and magnetic fields. One of the most interesting phenomena observed is the structuring of the barium cloud: the "fingering" or "striating" of the backside of the cloud. It is generally believed that this structuring is produced by the $\underline{E} \times \underline{B}$ instability [Linson and Workman, 1970]. A substantial research effort, both theoretical and computational, has been devoted to understanding this instability and its application to barium cloud structuring. The purpose of this paper is to examine the linear stability of the steepened backside of two-dimensional barium clouds. We derive a relatively simple expression for the growth rate of the perturbations on these backsides which are functions of (1) the parameter M which is the ratio of the integrated Pedersen conductivity inside the barium cloud to that of the background ionosphere; (2) the shape of the cloud; (3) the plasma density gradient scale length L of the steepened backside; and (4) the perturbation wavenumber k. For sufficiently small kL, the growth time as a function of M traces out the "U-shaped curve" observed in previous studies [Linson, 1975; McDonald et al., 1981; Overman and Zabusky, 1980].

The outline for the remainder of this paper is as follows. In Section II we derive the growth rate of a two-dimensional barium cloud, where we approximate the steepened backside of the cloud as a discontinuity. In Section III we relax this constraint and let the density gradient scale length on the steepened backside be finite. In Section IV we relate our results to the "U-shaped curve" of critical Reynold's number seen by McDonald et al. [1981]. Finally, in Section V we present our conclusions and plans for future work.

Manuscript approved February 9, 1984.

II. Theory

We are interested in deriving the approximate linear growth rate for perturbations applied to the steepened backside of an ionospheric barium cloud. To achieve this we need several earlier results. The first is the equation derived by Ossakow and Chaturvedi [1978] which describes the rise velocities of plasma depletions in the shape of circular or ellipsoidal cylinders in the presence of a gravitational field perpendicular to a magnetic field $\underline{B}$. It is simple matter to extend their equation to plasma enhancements in the presence of an electric field $\underline{E}_o$ (i.e., barium clouds). The second result we shall need is the equation for the linear growth rate of the gradient drift $\underline{E} \times \underline{B}$ instability when the plasma profile consists of a single discontinuity in plasma density. This result has been derived by Huba and Zalesak [1983]. Finally, we shall need to know the dependence of the growth rate on the plasma density gradient scale length in the real-world regime between the two extremes of a density discontinuity on the one hand, and a smooth density distribution, where "local" theory should hold, on the other. For this information we draw upon the work of Huba et al. [1982] and a suggestion of J.A. Fedder [private communication, 1983].

Consider an idealized two-dimensional barium cloud consisting of a uniform density inside of a "waterbag" of ellipsoidal shape, acted on by a uniform external electric field $E_o \hat{x}$ in the presence of a magnetic field $B \hat{z}$. Let the ellipse axis in the direction of $\underline{E}_o$ be denoted by b and the other axis by a. Let the ion density of the barium cloud be denoted by n_b and that of the background ionosphere be denoted by n_o. This situation is depicted in Fig. 1. Then the results of Ossakow and Chaturvedi [1978] can be extended to show that the total electric field $\underline{E}$ inside the barium cloud is given by

$$\underline{E} = \underline{E}_o \frac{1 + R}{1 + MR} \qquad (1)$$

where

$$R \equiv a/b \qquad (2)$$

$$M \equiv (n_b + n_o)/n_o \qquad (3)$$

Equation (1) is the first result we shall need for our analysis.

Now consider a one-dimensional barium cloud, infinite and invariant in the x direction, which consists of a single jump in electron density from $n_b + n_o$ to n_o at $y = 0$. That is, $n_e = n_b + n_o$ for $y < 0$, and $n_e = n_o$ for $y > 0$, where n_e is electron density. Again we impose an electric field, denoted here by $\underline{E}$, aligned in the x direction. We then perturb the interface at $y = 0$ such that the position of the perturbed interface $\eta(x)$ is given by $\eta = \alpha \cos kx$ where α is a small distance. This situation is depicted in Fig. 2. It is shown in Huba and Zalesak [1983] that this interface will grow in amplitude with growth rate γ given by

$$\gamma = k \frac{cE}{B} \frac{M - 1}{M + 1} \qquad (4)$$

Equation (4) was derived under the assumption that both Hall currents and the inertial terms in the ion momentum equation are negligible. M is as defined in Eq. (3). An alternative derivation of Eq. (4), due to the first author, is given in the Appendix.

Consider a steepened barium cloud, as depicted in Fig. 3. If we consider only the tip of the steepened backside, and we further consider only perturbations of sufficiently long wavelength such that kL is small, where k is the perturbation wavenumber and L is the characteristic scale length of the density gradient on the backside of the cloud, then the growth rate of these perturbations should be well approximated by Eq. (4). However, note that the electric field $\underline{E}$ seen by the cloud is not the externally imposed electric field $\underline{E}_o$, but rather some smaller electric field that is a result of the shielding effects of the two-dimensional cloud. If we approximate our barium cloud as an ellipse as depicted in Fig. 1, then Eq. (1) yields precisely this smaller electric field:

$$\underline{E} = \underline{E}_o \frac{1 + R}{1 + MR} \qquad (5)$$

Substituting this field into Eq. (4) we get finally

$$\gamma = k \frac{cE_o}{B} \frac{(M - 1)\ (R + 1)}{(M + 1)\ (MR + 1)} \tag{6}$$

For the special case of a circular shape (R = 1) we get

$$\gamma = 2k \frac{cE_o}{B} \frac{M - 1}{(M + 1)^2} \tag{7}$$

Equation (7) traces out an inverted "U-shaped" dependence of γ on M, which we relate to the "U-shaped curve" of McDonald et al. [1981] in Section IV. The M dependence of γ is given by $(M - 1)/(M + 1)^2$, which maximizes at M = 3. A plot of this quantity, as well as its two factors, $(M - 1)/(M + 1)$ and $(M + 1)^{-1}$, which are due to 1-D instability and 2-D shielding respectively, is shown in Fig. 4. Note that the falloff in γ for small M is due to 1-D instability effects alone, while the falloff in γ for large M is due to 2-D shielding alone. Thus we conclude that the resistance of 2-D high M clouds to bifurcation (see McDonald et al., 1981) is due to 2-D shielding.

Equation (7) is very similar, but not identical to, an expression derived by Overman and Zabusky [1980] for the case of a perfectly circular "waterbag" model of a barium cloud. Their linearized analysis showed the results of perturbing a circular cloud of radius r_o by making radial sinusoidal displacements of its boundary

$$r = r_o + \sum_m \alpha_m \cos m\phi \tag{8}$$

where r_o is the initial radius of the cloud, r is the perturbed radius, α_m is the perturbation amplitude, ϕ is the angle measured from the backside of the cloud, and m is an integer (mode number). They found that

$$\partial\ \alpha_m/\partial t = \left(2 \frac{m}{r_o} \frac{cE_o}{B} \frac{M - 1}{(M + 1)^2}\right) \alpha_{m+1} \tag{9}$$

Note that if we write our perturbation (8) in terms of arc length $r_o\phi$ multiplied by a wavenumber k then $\cos m\phi = \cos k\, r_o\phi$ and

$$m/r_o = k \tag{10}$$

making Eq. (9) yield the same "growth rate" as given by Eq. (7), except that there is an α_m on the left side of (9) and an α_{m+1} on the right side. Thus we have not totally recovered the results of Overman and Zabusky [1980] for the circular waterbag, but the similarity is both striking and encouraging. Furthermore, our formula Eq. (6) exhibits several advantages over the Overman and Zabusky [1980] result. For one thing, we do not have the ambiguity associated with using the term "growth rate" to describe the rate at which mode m+1 supplies energy to mode m (see Eq. 9). Equation (6) is a genuine growth rate in the usual sense of the word. Also, precisely because we have derived Eq. (6) as the product of a term involving global effects, i.e., the overall shape of the cloud, and a term involving local effects, i.e., the densities and electric fields at the steepened edge of the cloud, it is possible to generalize the results to elliptical clouds (via the R dependence) and, more importantly, to the case of a finite density gradient on the backside of the cloud, as we shall see in Section III.

III. Approximating the Growth Rate in Cases Where the Density Gradient is Finite

Let us return to Eq. (4). This equation was derived under the assumption that the edge of the barium cloud represented a true discontinuity in electron density; but real barium clouds consist of finite gradients. How accurately does Eq. (4) represent the case of a sharp but finite density gradient and how are we to treat the case where the density gradient is not even sharp? The work of Huba et al. [1983] gives numerically generated growth rates for a wide range of ratios of perturbation wavelength to gradient scale length. However, we are interested here in finding a simple, closed form expression. We propose the following, after a suggestion due to J.A. Fedder [private communication, 1983]

$$\gamma = k \frac{cE}{B} \frac{M-1}{M+1}, \quad kL < \frac{M+1}{M-1} \tag{11a}$$

$$\gamma = \frac{cE}{B} L^{-1}, \quad kL > \frac{M+1}{M-1} \tag{11b}$$

where L^{-1} is the maximum value of the quantity $\frac{1}{n_e}\frac{\partial n_e}{\partial y}$ for the geometry shown in Fig. 2 wherein we have replaced the discontinuity by a finite transition length. Equation (11b) is the classical result for the gradient drift instability from local theory (see Linson and Workman [1970]). The combined Eq. (11) agrees very well with the actual numerical results shown in Huba et al. [1983] when allowance is made for the fact that the L used in that paper differs from the one used here by a factor of 1.45. A comparison is shown in Fig. 5.

Combining Eq. (11) with Eq. (1), we get our final expressions for the growth rate of instabilities on the steepened backside of a barium cloud, in the absence of diffusion:

$$\gamma = k\,\frac{cE_o}{B}\,\frac{M-1}{M+1}\,\frac{R+1}{MR+1}\;, \quad kL < \frac{M+1}{M-1} \tag{12a}$$

$$\gamma = \frac{cE_o}{B}\,L^{-1}\,\frac{R+1}{MR+1}\;, \quad kL > \frac{M+1}{M-1} \tag{12b}$$

An obvious caveat to the above is that the perturbation wavelength $\lambda = k/2\pi$ must be smaller than the total diameter of the barium cloud to have the concept of a Fourier mode perturbation to make any sense at all.

We note that we have constructed our final expressions such that they not only very closely recover the results of Overman and Zabusky [1980] (for $R = 1$ and $L = 0$, as previously discussed) but also accurately reproduce the other known case of a one-dimensional plasma distribution ($R = 0$) for the entire range of kL.

IV. Relationship to the U-Shaped Curve

Equation (7), when viewed as an expression for γ as a function of M, yields a curve which maximizes at $M = 3$ and monotonically decreases away from this value for all values of M for which our barium cloud represents a true enhancement in electron density ($1 < M < \infty$). In fact the growth rate γ is zero for $M = 1$ and $M = \infty$. The corresponding curve of a growth time γ^{-1} versus M minimizes at $M = 3$ and goes to infinity at $M = 1$ and $M = \infty$. Curves of γ^{-1} versus λ which exhibit this qualitative behavior have been dubbed "U-shaped curves" by several researchers. If one views barium cloud bifurcation as a battle between the gradient drift growth rate γ and diffusive processes, then this U-shaped curve can be related to a similar U-shaped curve which relates the amount of diffusion D needed to stop the bifurcation of a cloud of a given radius r_o, as a function of M. Let us assume that the effect of diffusion is to contribute a term $-k^2D$ to the growth rate. Let us further approximate our barium cloud as a cylinder (R = 1), with very steep edges ($L/r_o < 1$). Then our growth rate is

$$\gamma = 2k \frac{cE_o}{B} \frac{M - 1}{(M + 1)^2} - k^2D \tag{13}$$

We will further assume that bifurcation, if it takes place at all, takes place at $kL \sim 1$. The bifurcation condition is then found by setting $\gamma = 0$ in (13) with $k = L^{-1}$:

$$L^{-2}D = 2L^{-1} \frac{cE_o}{B} \frac{M - 1}{(M + 1)^2} \tag{14}$$

$$\frac{\frac{cE_o}{B} L}{D} = \frac{1}{2} \frac{(M + 1)^2}{(M - 1)} \tag{15}$$

Noting that the left-hand side of (15) is simply McDonald et al.'s [1981] Reynolds number Re we get

$$\text{Re} = \frac{1}{2} \frac{(M + 1)^2}{(M - 1)} \tag{16}$$

Note that McDonald et al. [1981] got

$$\text{Re} \sim 75 \frac{(M + 1)^2}{(M - 1)} \tag{17}$$

so the functional forms are the same but our Eq. (16) seems to state that it requires 150 times more diffusion to stop bifurcation then was observed by McDonald et al. [1981]. Whence the discrepancy? Part of the difference can be removed by noting that the "finger" of McDonald et al. [1981] corresponds to taking $R = \infty$, in which case we would get

$$\text{Re} = \frac{M(M + 1)}{(M - 1)} \tag{18}$$

For large M this removes half of the discrepancy. Accounting for a larger share of the difference is McDonald et al.'s [1981] equation for computing L (their Eq. (5)), which appears to give L's which are larger than the minimum gradient scale length we have assumed here. However, we do not believe this effect could account for a factor greater than ten, which still leaves a discrepancy of more than seven. That is, the model we propose here predicts a factor of at least seven greater diffusion necessary to stop bifurcation than is given by McDonald et al. [1981]. The U-shaped dependence on M, however, is the same. We hope to be able to resolve these quantitative differences in the sequel to this paper [Zalesak, 1983], where the full effects of diffusion are considered.

V. Conclusions

We have derived an expression, Eq. 12, that we believe to be a good approximation to the actual growth rate of perturbations applied to the steepened backside of a two-dimensional ionospheric barium cloud, for the case of no electron or ion diffusion. Diffusion plays a large role in the behavior of two-dimensional barium clouds, however, and affects the growth rate in two ways:

(i) diffusion will determine the maximum density gradient that can exist at the backside of the barium cloud;

(ii) diffusion will provide a damping mechanism for any perturbation which attempts to grow on the backside of the cloud. We have already treated this effect, at least approximately, in Section IV.

These two combined effects of diffusion are treated in the sequel to this paper [Zalesak, 1983].

Appendix. Derivation of the $\underline{E} \times \underline{B}$ Instability Growth Rate for a Discontinuous Interface

The geometry we shall consider is depicted in Fig. 6. The differential equations describing this system are [McDonald et al., 1981]:

$$\frac{\partial n_e}{\partial t} + \nabla \cdot (n_e \underline{v}) = 0 \tag{A1}$$

$$\nabla \cdot (n_e \nabla \phi) = \underline{E} \cdot \nabla n_e \tag{A2}$$

where

$$\underline{v} = -\frac{c}{B^2} \nabla \phi \times \underline{B} \tag{A3}$$

Here n_e is the electron density, which is n' for $y > \eta(x)$ and n for $y < \eta(x)$, c is the speed of light, $\underline{B}$ is the magnetic field, $\underline{E}$ is the imposed electric field, and ϕ is the induced electrostatic potential. Note from Eq. (A3) that we have placed ourselves in a reference frame drifting with the $\underline{E} \times \underline{B}$ velocity. Our two-dimensional plasma resides in the x-y Cartesian plane. $\underline{B}$ is aligned along $\hat{z}$.

If $\underline{E} = E\,\hat{x}$, and we ignore Hall currents, then the geometry depicted in Fig. 6 is an equilibrium configuration with $\eta = 0$. We now apply a perturbation to the position of the interface, denoted by $\eta(x)$, of the form

$$\eta = \alpha \cos kx\, e^{i\sigma t} \tag{A4}$$

and ask what the response of the electrostatic potentials in the two regions will be. Noting that $\nabla n = \nabla n' = 0$, since the fluid is incompressible and we have only moved the interface, Eq. (A2) becomes

$$\nabla^2 \phi' = 0 \quad , \quad y > \eta(x) \tag{A5}$$

$$\nabla^2 \phi = 0 \quad , \quad y < \eta(x) \tag{A6}$$

If we assume solutions of the form

$$\phi' = G'(y) \sin kx\, e^{i\sigma t} \quad , \qquad y > \eta(x) \tag{A7}$$

$$\phi = G(y) \sin kx\, e^{i\sigma t} \quad , \quad y < \eta(x) \tag{A8}$$

(Note we are <u>assuming</u> a specified phase relationship between η and ϕ which will later be shown to satisfy our boundary conditions), then demanding that ϕ' vanish at $y = +\infty$ and ϕ vanish at $y = -\infty$ yields

$$\phi' = C' e^{-ky} \sin kx\, e^{i\sigma t}, \quad y > \eta(x) \tag{A9}$$

$$\phi = C \; e^{+ky} \sin kx\, e^{i\sigma t}, \quad y < \eta(x) \tag{A10}$$

We shall later see that σ is pure imaginary, meaning that the perturbation does not <u>propagate</u> in the x direction. Therefore the movement of the interface is determined by the y-components of the velocity fields at the interface, <u>which must be the same</u> for the interface to be well defined (no interpenetration or cavitation):

$$\frac{\partial \eta}{\partial t} = v_y' = v_y \text{ at } y = \eta \tag{A11}$$

but

$$v_y = \frac{c}{B}\frac{\partial \phi}{\partial x}$$

$$v_y' = \frac{c}{B}\frac{\partial \phi'}{\partial x}$$

so

$$\frac{B}{c}\frac{\partial \eta}{\partial t} = kC' e^{-ky} \cos kx\, e^{i\sigma t}\Big|_{y = \eta} \tag{A12}$$

$$= kC\, e^{+ky} \cos kx\, e^{i\sigma t}\Big|_{y = \eta}$$

Since η is assumed to be arbitrarily small we have

$$C' = C \tag{A13}$$

but by Eq. (A4)

$$\frac{\partial \eta}{\partial t} = i\sigma\, \alpha \cos kx\, e^{i\sigma t} \tag{A14}$$

so comparing (A12) and (A14) we get

$$i \frac{B}{c} \sigma \alpha = kC' = kC \tag{A15}$$

Now we have one last constraint to satisfy, namely that $\nabla \cdot \underline{J} = 0$. We will use the integral form of this constraint, that is

$$\iint_{S(V)} \underline{J} \cdot \hat{n} \, dS = 0 \tag{A16}$$

for any control volume V bounded by surface S(V). Equation (A16) obviously holds for any control volume not containing the interface $\eta(y)$, since n = a constant and $\nabla^2 \phi = 0$. Consider our interface and the control volume depicted in Fig. 7. J_L, J_R, J_B and J_T are the total currents flowing through the left, right, bottom and top surfaces respectively of the control volume. We define $\eta_R \equiv \eta(x_R)$ and $\eta_L \equiv \eta(x_L)$.

$$\iint \underline{J} \cdot \hat{n} \, ds = J_R - J_L + J_T - J_B = 0 \tag{A17}$$

$$J_T = - n \frac{\partial \phi'}{\partial y} = + n' kCe^{-ky_T} \int_{x_L}^{x_R} \sin kx \, dx \tag{A18}$$

$$J_B = - n \frac{\partial \phi}{\partial y} = - nkCe^{ky_B} \int_{x_L}^{x_R} \sin kx \, dx \tag{A19}$$

$$J_R = J_R^I + J_R^{II} \tag{A20}$$

where

$$J_R^{II} \equiv \int_{y_B}^{\eta_R} n\left(E - \frac{\partial \phi}{\partial x}\right) dy \tag{A21}$$

$$J_R^I \equiv \int_{\eta_R}^{y_T} n'\left(E - \frac{\partial \phi'}{\partial x}\right) dy \tag{A22}$$

Since both $\frac{\partial \phi'}{\partial x}$ and the integration path length $(y_T - \eta_R)$ are perturbation quantities in Eq. (A22), the integral involving $n' \, \partial \phi'/\partial x$ will be quadratic in the perturbation and hence negligible. We then get

$$J_R^I = \int_{\eta_R}^{y_T} n' E \, dy = (y_T - \eta_R) n' E \tag{A23}$$

Similarly

$$J_R^{II} = \int_{y_B}^{\eta_R} n\ E\ dy = (\eta_R - y_B)nE \tag{A24}$$

$$J_R = J_R^I + J_R^{II} = y_T n'E - y_B nE + \eta_R(n - n')E \tag{A25}$$

Using similar reasoning we get

$$J_L = y_T n'E - y_B nE + \eta_L(n - n')E \tag{A26}$$

$$J_R - J_L = (n - n')E(\eta_R - \eta_L) \tag{A27}$$

We now let

$$(x_R - x_L) \to 0 \tag{A28}$$

Then

$$J_R - J_L = (n - n')E \frac{\partial\eta}{\partial x} (x_R - x_L) \tag{A29}$$

but from Eq. (A4)

$$\frac{\partial\eta}{\partial x} = -\ \alpha k \sin kx\ e^{i\sigma t} \tag{A30}$$

so

$$J_R - J_L = -\ (n - n')E\ \alpha\ k \sin kx\ e^{i\sigma t}(x_R - x_L). \tag{A31}$$

Using Eq. (A28) and noting that y_T and y_B are perturbation quantities we can set

$$\int_{x_L}^{x_R} \sin kx\ dx \approx \sin kx(x_R - x_L) \tag{A32}$$

$$e^{-ky_T} = e^{+ky_B} \approx 1 \tag{A33}$$

and hence Eqs. (A18) and (A19) become

$$J_T = + n' kC \sin kx(x_R - x_L) \tag{A34}$$

$$J_B = - n kC \sin kx(x_R - x_L) \tag{A35}$$

$$J_T - J_B = (n' + n) kC \sin kx(x_R - x_L) \tag{A36}$$

Combining Eq. (A36), (A31), and (A17) and canceling the common term $k \sin kx\, e^{i\sigma t}(x_R - x_L)$ we get

$$(n' + n)C - (n - n')E \alpha = 0 \tag{A37}$$

Now substituting Eq. (A15) into (A37)

$$(n' + n)i \frac{B\sigma\alpha}{kc} = (n - n')E\alpha$$

$$\sigma = - i \frac{n - n}{n + n} k \frac{cE}{B} \tag{A38}$$

$$\gamma = \frac{n - n}{n + n'} k \frac{cE}{B} \tag{A39}$$

Defining $M \equiv n/n'$

$$\gamma = \frac{M - 1}{M + 1} k \frac{cE}{B} \tag{A40}$$

which corresponds to the result obtained by Huba and Zalesak [1983].

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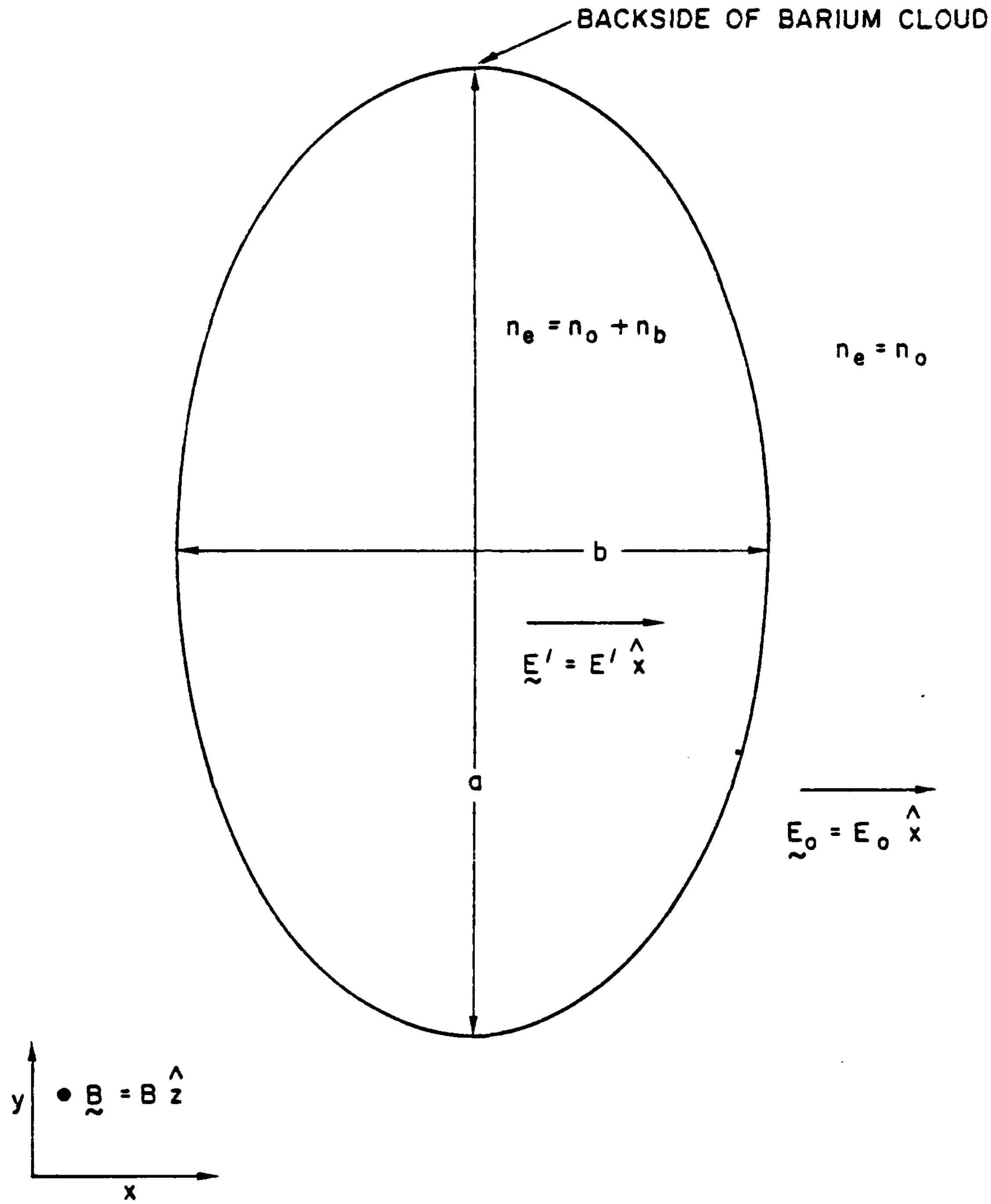

Fig. 1. Idealized "waterbag" elliptical barium cloud geometry. The total electron density, barium cloud ion density, and background ionosphere ion density are denoted by n_e, n_b, and n_o respectively. $\underline{B}$ is the ambient magnetic field, and $\underline{E_o}$ is the externally imposed electric field. $\underline{E}$ is the constant shielded electric field inside the cloud. a and b are the major and minor axes of the ellipse respectively.

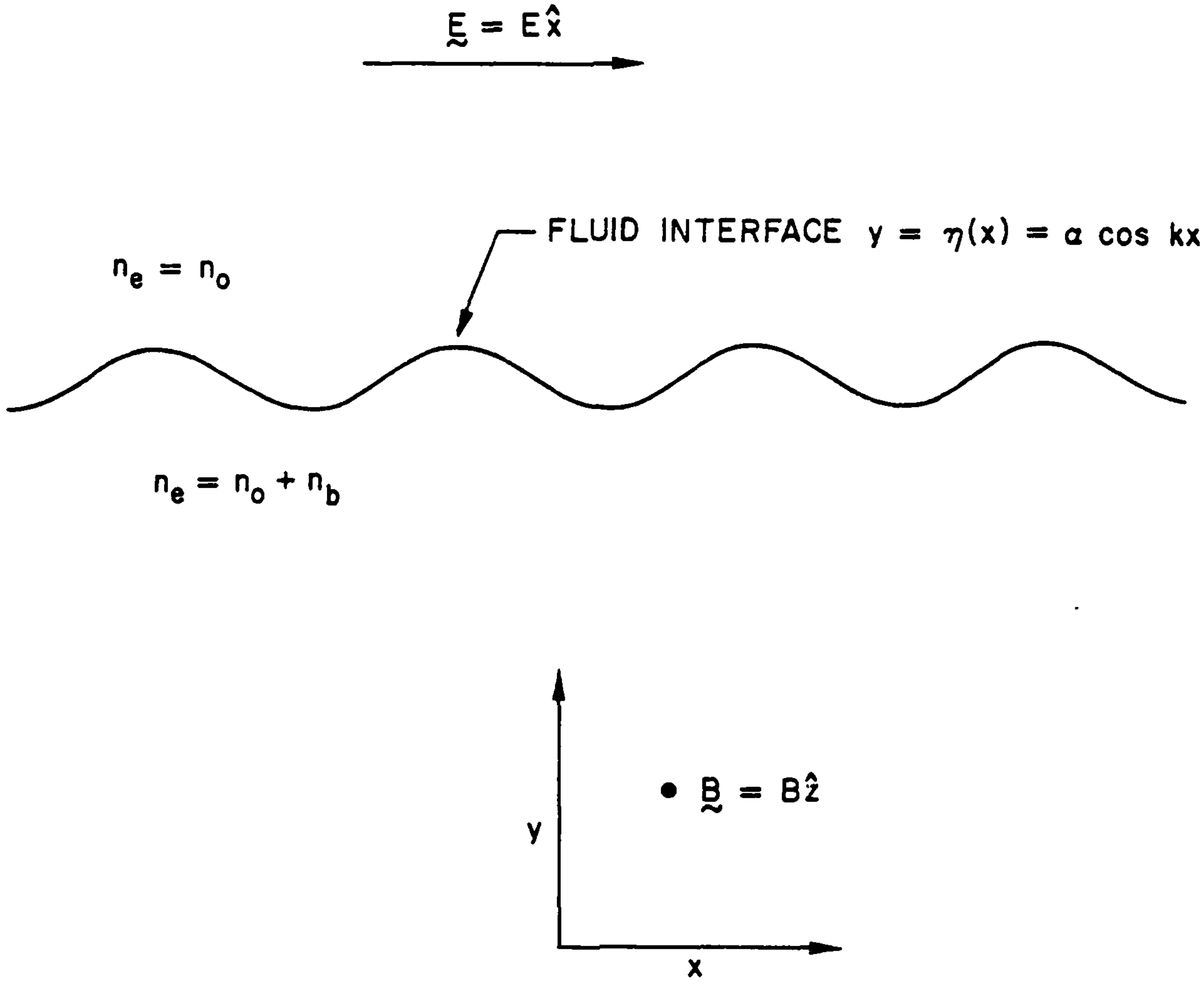

Fig. 2 Geometry of 1-D barium cloud configuration for stability analysis of discontinuous interface.

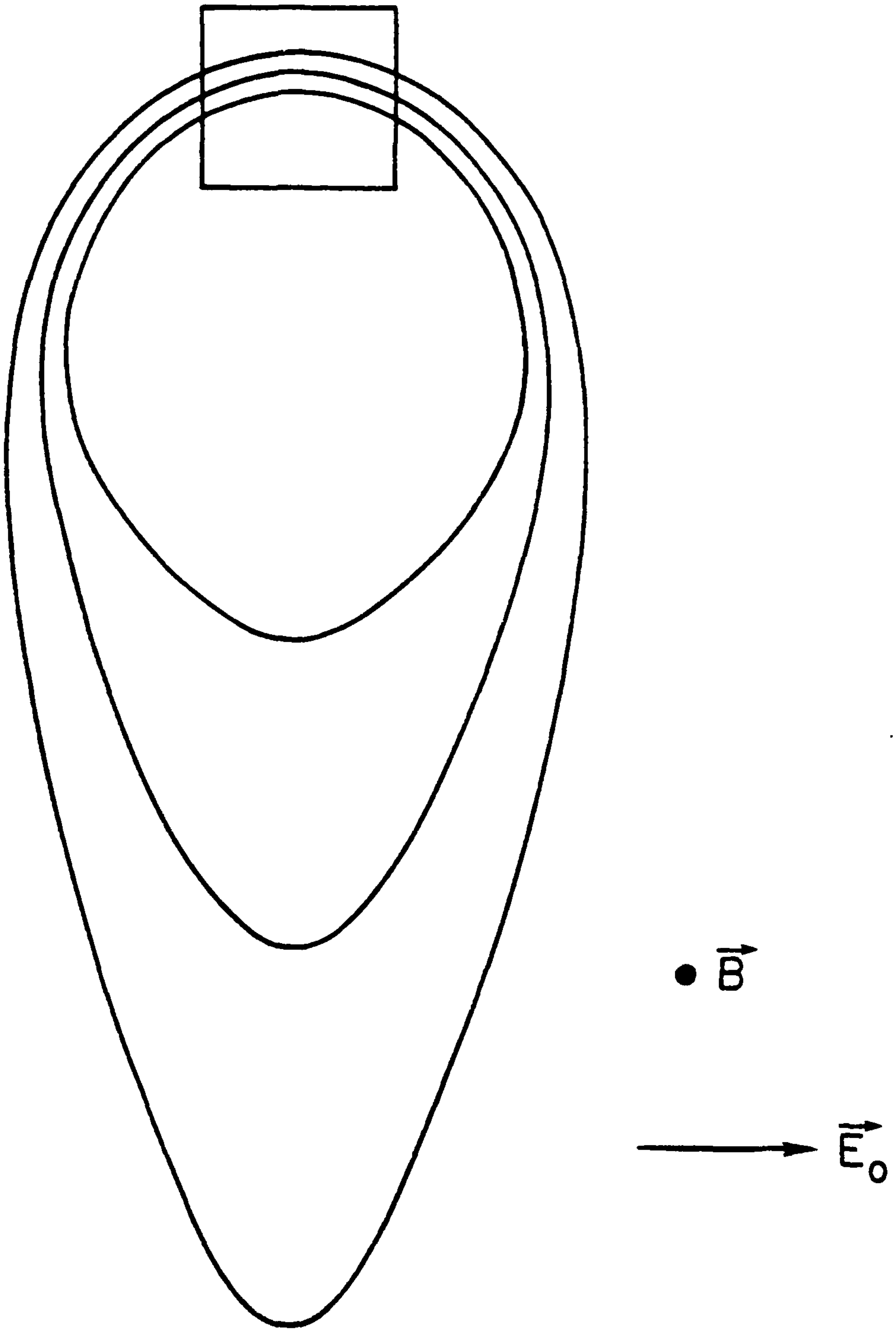

Fig. 3 Isodensity contours of a steepened 2-D barium cloud. The box indicates that region of space over which we will apply our 1-D stability analysis by approximating the density gradient there as a 1-D, slowly evolving structure.

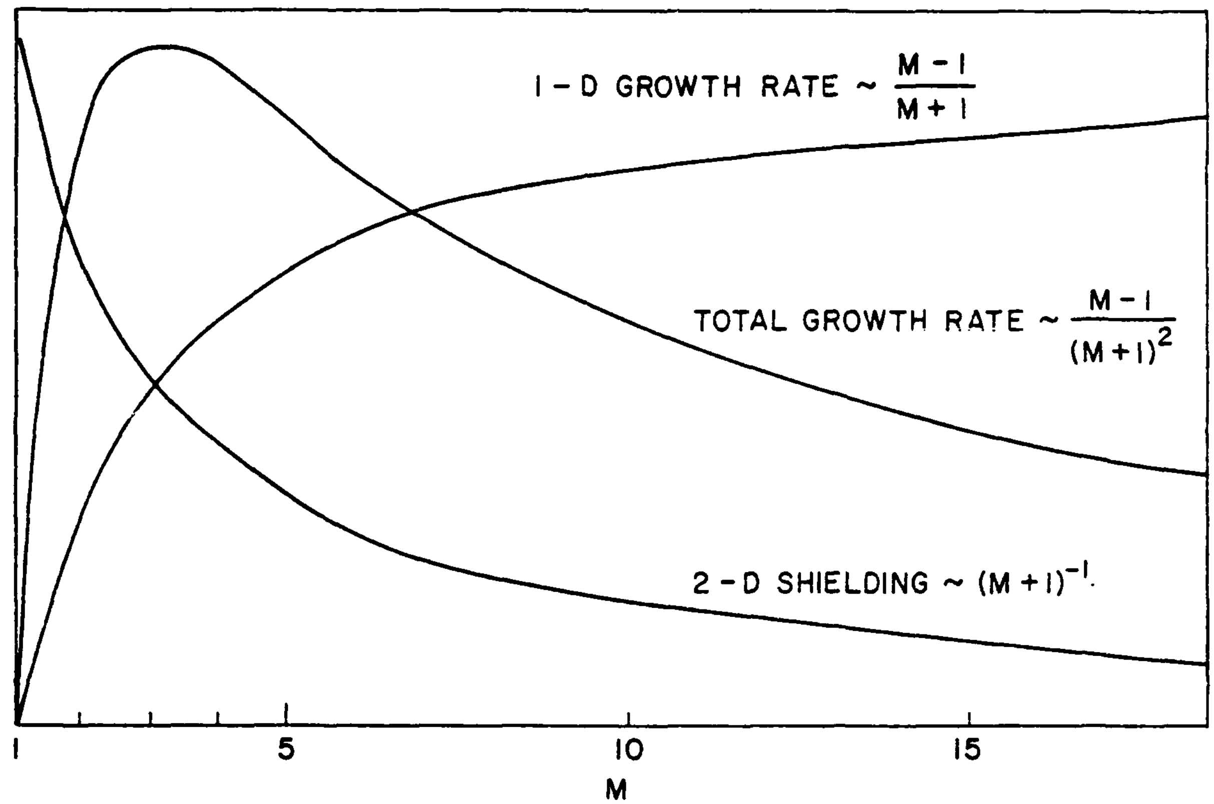

Fig. 4 Plot of the M-dependence of the growth rate given by Eq. (7), $M - 1/(M + 1)^2$. Also shown are its two component parts, $(M - 1)/(M + 1)$ and $(M + 1)^{-1}$, due to 1-D instability and 2-D shielding results respectively.

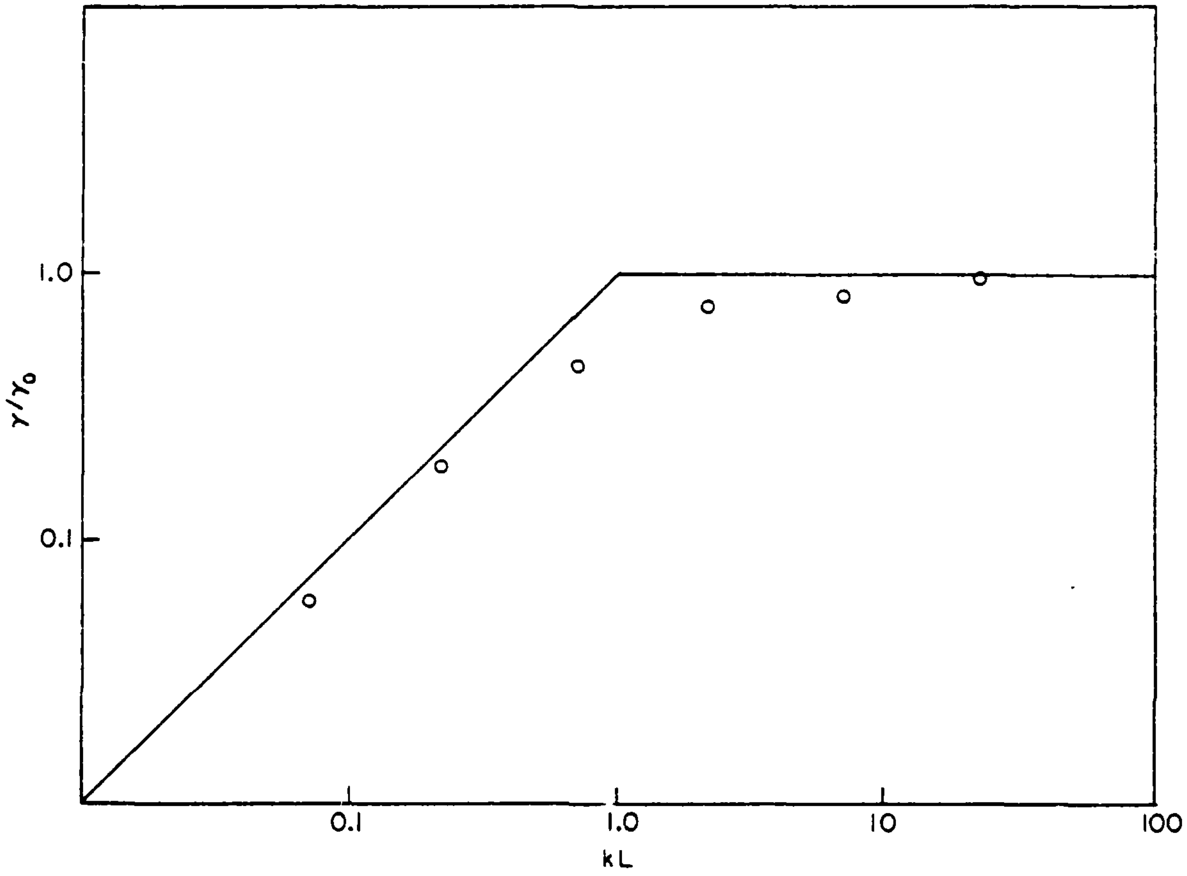

Fig. 5 Plot of normalized growth rate versus kL for the case of the M = 39 cloud considered by Huba et al. [1982]. The solid line depicts the simple model shown here, while the circles represent numerical results from Huba et al. [1982].

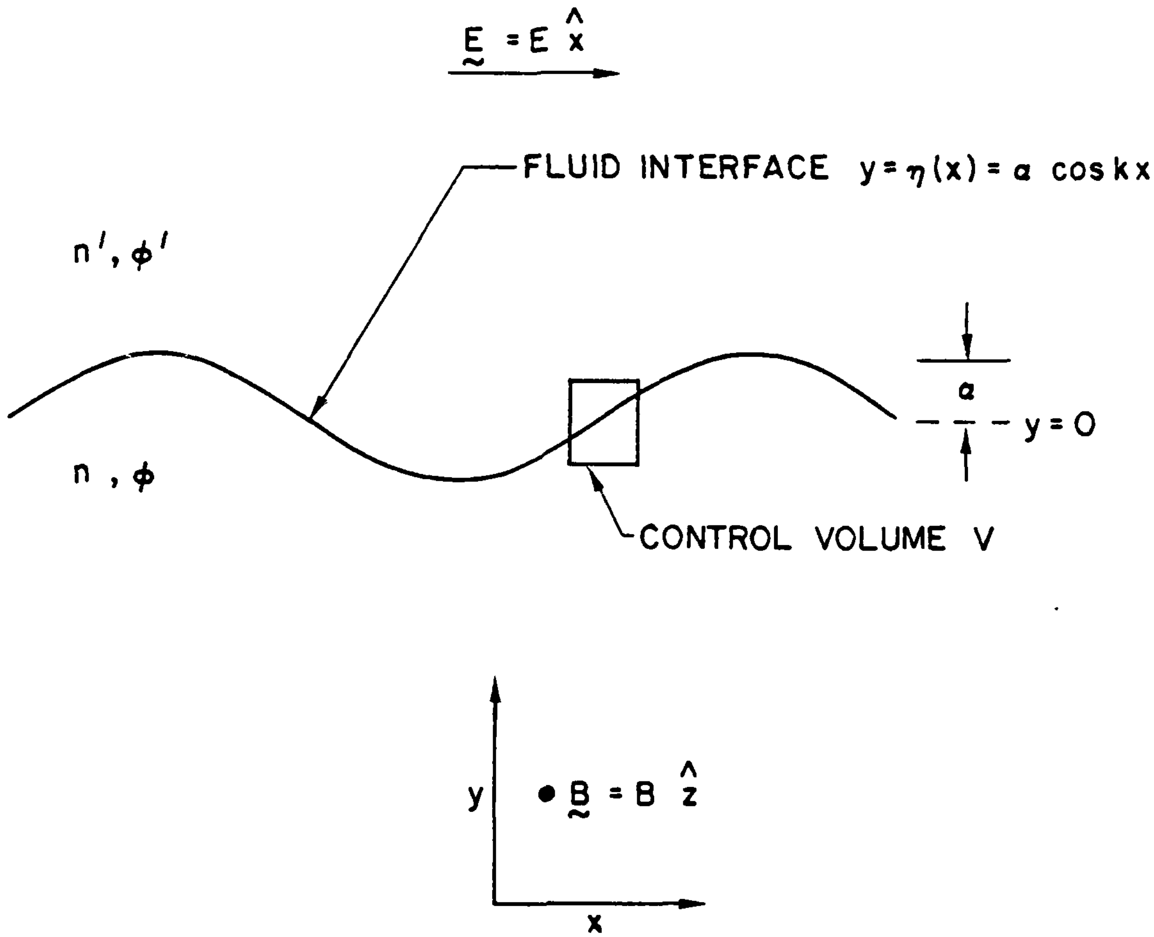

Fig. 6 Geometry of 1-D barium cloud configuration for stability analysis, as in Fig. 2. n and ϕ are the plasma density and electrostatic potential, respectively. Primed and unprimed variables denote quantities above and below the interface $y = \eta(x)$, respectively. Shown is the control volume used to evaluate $\nabla \cdot J$.

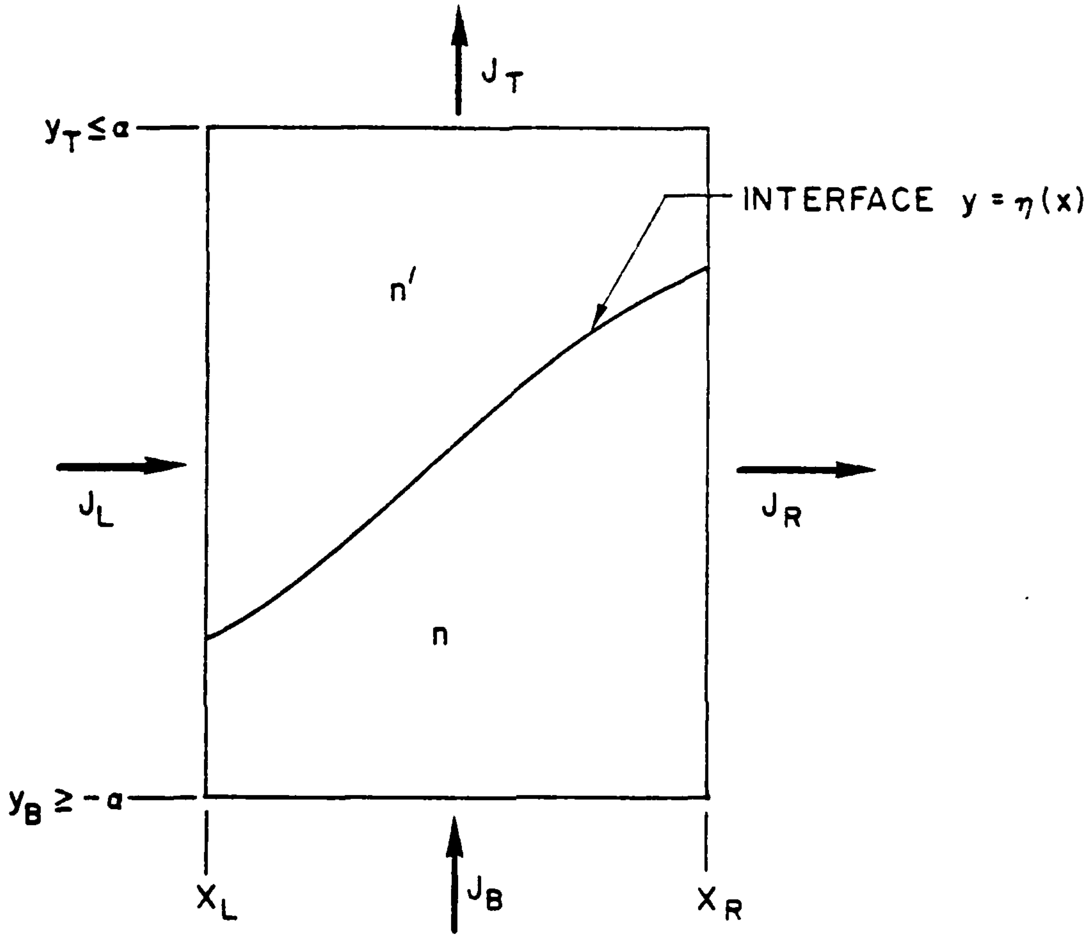

Fig. 7. Blow-up of the control volume shown in Fig. 6. Depicted are the total currents flowing through the four surfaces of the volume.

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